

The spectrum of scalar-meson nonets in the Resonance-Spectrum Expansion*

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Abstract

We argue that the low-lying scalar-meson nonet [1] makes part of a subset of a family of infinitely many scalar-meson nonets, which in turn makes part of a family of infinitely many quark-antiquark bound states and resonances. We outline the properties of this subset.

1 Introduction

Except for a few mesons, like pions and kaons, most quark-antiquark states show up as resonances in systems of two or more mesons. It is thus opportune to study the interplay of meson-meson scattering and $q\bar{q}$ confinement [2, 3].

The mesonic resonances extracted from experiment are organized by flavor content, $J^{PC}I^G$ quantum numbers, mass, and width. From the few hundred listed in the PDG tables [4], one would not yet conclude that they are abundant. Nevertheless, based on the $b\bar{b}$ and $c\bar{c}$ spectra, we concluded in Ref. [2] that there must exist an infinity of such states, though cut off from observation at higher masses because of the many two-meson systems coupling to $q\bar{q}$. Accordingly, we expect an infinity of scattering poles to show up in meson-meson scattering, here represented by

$$E = P_0, P_1, P_2, \dots \quad (1)$$

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Unitarity then requires that in the one-channel restriction, assuming the poles (1) to be simple poles, the elastic scattering matrix S be given by¹

$$S(E) = \frac{(E - P_0^*)(E - P_1^*)(E - P_2^*) \dots}{(E - P_0)(E - P_1)(E - P_2) \dots} . \quad (2)$$

If we assume that the resonances (1) stem from an underlying confinement spectrum, given by the real quantities

$$E = E_0, E_1, E_2, \dots , \quad (3)$$

then we may represent the differences $(P_n - E_n)$, for $n = 0, 1, 2, \dots$, by ΔE_n . Thus, we obtain for the unitary S -matrix the expression

$$S(E) = \frac{(E - E_0 - \Delta E_0^*)(E - E_1 - \Delta E_1^*)(E - E_2 - \Delta E_2^*) \dots}{(E - E_0 - \Delta E_0)(E - E_1 - \Delta E_1)(E - E_2 - \Delta E_2) \dots} . \quad (4)$$

So we assume here that resonances occur in scattering because the two-meson system couples to confined states, usually of the $q\bar{q}$ type, viz. in non-exotic meson-meson scattering. Let the strength of the coupling be given by λ . For vanishing λ , we presume that the widths and real shifts of the resonances also vanish. Consequently, the scattering poles end up at the positions of the confinement spectrum (3), and so

$$\lim_{\lambda \downarrow 0} \Delta E_n = 0 \quad \text{for } n = 0, 1, 2, \dots . \quad (5)$$

As a result, the scattering matrix tends to unity, as expected in case there is no interaction.

An obvious candidate for an expression of the form (4) looks like

$$S(E) = \left[1 + \lambda^2 \left\{ \sum_n \frac{G(E)^*}{E - E_n} \right\} \right] \left[1 + \lambda^2 \left\{ \sum_n \frac{G(E)}{E - E_n} \right\} \right]^{-1} , \quad (6)$$

where G is a smooth complex function of energy E .

¹ Note that we do not consider here a possible overall phase factor representing a background.

2 Kaon-pion S-wave scattering

In order to compare expression (6) with results of experiment, we must choose a suitable complex function G . This has been done in Ref. [2], and was further developed in Refs. [5–8]. Furthermore, values for the real spectrum (3) must be chosen. In principle, one could fit E_n ($n = 0, 1, 2, \dots$) to experiment. But it is our experience that the spectrum listed in Ref. [4] is not yet rich enough to determine a suitable confinement spectrum. In Refs. [2, 5–8] we proposed flavor masses and a universal level spacing ω for this purpose. The latter quantity can, with some confidence, be deduced from the $c\bar{c}$ and $b\bar{b}$ spectra, and also from the light positive-parity resonances [9]. One finds $\omega = 0.19$ GeV, corresponding to interquark distances ranging from 0.2 fm for $b\bar{b}$ to 0.6 fm for light quarks..

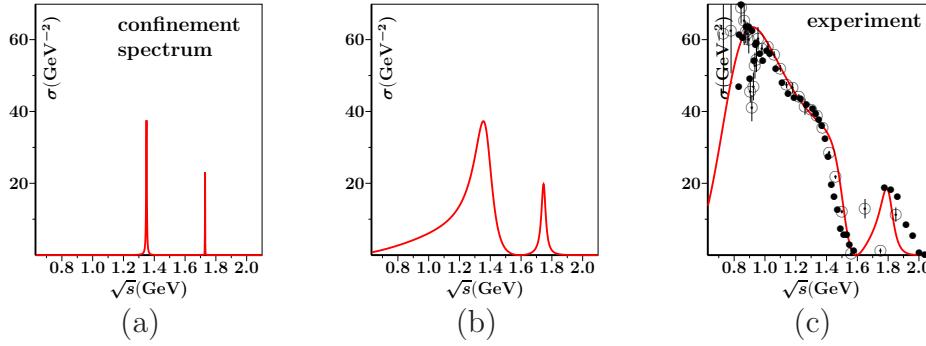


Figure 1: Cross section for S -wave isodoublet $K\pi$ scattering. Left: for very small values of λ , one observes the $J^P = 0^+$ $n\bar{s}$ confinement spectrum. Middle: when λ takes about half its model value, one notices some more structure for low invariant masses. Right: at the model value of λ , the latter structure becomes dominant and well in agreement with the experimental observations. The data are taken from Ref. [10] (open circles) and Ref. [11] (full circles).

After fitting the parameters to heavy-heavy, heavy-light and light-light vector and pseudoscalar data [5], we turn our attention to the scalar mesons [12]. This is just a matter of setting quantum numbers $J^P = 0^+$, determining E_0 , calculating cross sections from expression (6), and comparing to available data. Here, we will concentrate on the nonstrange-strange ($n\bar{s}$) centers of gravity of the scalar nonets, for energies up to about 2 GeV, and use the elastic-scattering data of Refs. [10, 11].

In Fig. 1 we show how cross sections following from formula (6) vary with increasing λ , for S -wave isodoublet $K\pi$ scattering. In Fig. 1a the $n\bar{s}$ confinement spectrum is well visible for small λ , whereas in Fig. 1c, for the model value of λ , experiment is reproduced. We find a fair agreement for total invariant masses up to 1.6 GeV.

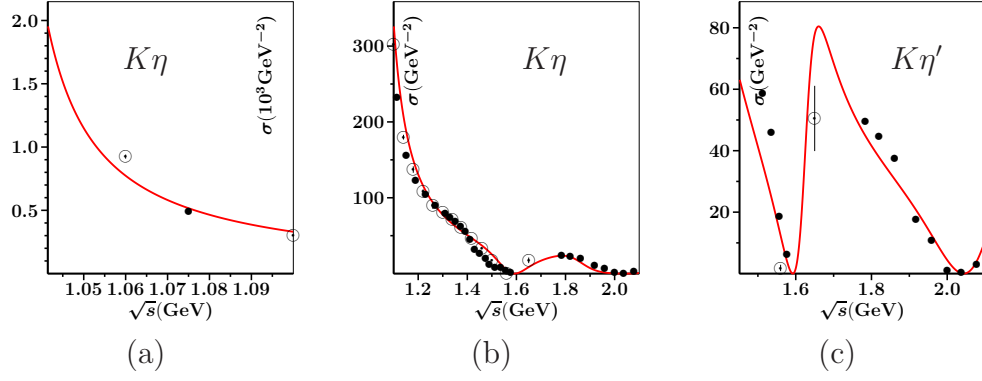


Figure 2: S -wave $K\eta$ and $K\eta'$ “cross section” (see text), as a function of total invariant mass. (a): $K\eta$ from threshold up to 1.1 GeV. (b): $K\eta$ from 1.1 GeV up to 2.1 GeV. (c): $K\eta'$ from threshold up to 2.1 GeV. The data are taken from Ref. [10] (open circles) and Ref. [11] (full circles).

Now, in order to have some idea about the performance of formula (6) for S -wave $I = 1/2$ $K\pi$ scattering at higher energies, we argue that, as in our model there is only one non-trivial eigen-phase shift for the coupled $K\pi + K\eta + K\eta'$ system, we may compare the phase shifts of our model for $K\eta$ and $K\eta'$ to the experimental phase shifts for $K\pi$. We do this comparison in Fig. 2, where, instead of the phase shifts, we plot the cross sections, assuming no inelasticity in all cases. We observe an extremely good agreement. In particular, for $K\eta'$ (Fig. 2c) we become aware of a structure in the data at about 1.9 GeV, indicating the presence of a not anticipated pole. This is something we would not have easily noticed from the data alone.

3 A basketful of scalar nonets

When we inspect formula (6) for poles in the S -wave isodoublet $K\pi$ scattering amplitude, then we find the pole structure as summarized in Table 1, i.e., five poles at energies up to about 2.2 GeV real part. The first pole, at $0.772 - 0.281i$ GeV, describes the $K_0^*(800)$ structure [13], whereas the second pole, at $1.52 - 0.097i$ GeV, represents the well-established $K_0^*(1430)$ resonance [14].

Pole (GeV)	$0.772 - 0.281i$	$1.52 - 0.097i$	$1.79 - 0.052i$	$2.04 - 0.15i$	$2.14 - 0.065i$
Origin	continuum	confinement	confinement	continuum	confinement

Table 1: T -matrix poles for S -wave $K\pi$ scattering, as obtained from Eq. (6).

Our model is explicitly flavor independent, meaning that the only flavor breaking in formula (6) stems from the effective quark masses, which determine the ground state of the confinement spectrum (3), and from the masses of the mesons in the scattering channels. Consequently, $\pi\pi$ scattering is not very different from $K\pi$ scattering in our model. We may expect then that each of the two flavor combinations that couple to isoscalar S -wave $\pi\pi$ and KK scattering has a pole structure similar to the one in isodoublet $K\pi$ scattering, with the proviso that $n\bar{n}$ - $s\bar{s}$ mixing in the $I=0$ case introduces an extra complication [15]. Hence, since also $\eta\pi$ is similar to $K\pi$ in our model [16], with each pole of Table 1 we associate a full nonet of scalar mesons. The often read comment that *too many isoscalar states are observed* [17], in order to justify the application of alternative quark, or even quarkless, configurations [18], is not confirmed here.

4 Confinement and continuum poles

In order to explain the difference between *confinement* and *continuum* poles (see Table 1), we turn to another member of the scalar-meson family, namely the heavy-light ($c\bar{s}$) $D_{s0}^*(2317)$ meson [19].

The mass of the $D_{s0}^*(2317)$ ends up below the threshold of the lowest OZI-allowed decay mode (i.e., DK). Consequently, it represents a bound state in this specific selection of decay channels [20], which we consider the most important. Accordingly, the $D_{s0}^*(2317)$ may be represented by a pole on the real energy axis. The pole representing the first radial excitation of the $c\bar{s}$ system in a relative P -wave comes out well above the DK threshold. In Ref. [8], two poles were found, one at 2.32 GeV and a second at $(2.85 - i0.024)$ GeV, representing the ground state and the first radial excitation of the $J^P = 0^+$ $c\bar{s}$ system, respectively. Experiment [21] reported a $c\bar{s}$ structure at 2.86 GeV, with the same line-shape as our theoretical prediction [8], and being compatible with $J^P = 0^+$ quantum numbers.

But in Ref. [8] an additional pole showed up in the scattering amplitude. Its theoretical position was reported at $(2.78 - i0.23)$ GeV. In Fig. 3 we show the trajectories of the two lowest-

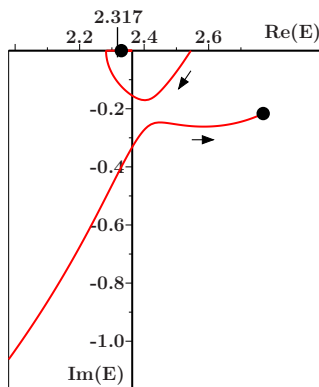


Figure 3: Trajectories of poles in the DK S -wave scattering amplitude for the model of Ref. [8], as a function of the amount of unquenching. In the quenched approximation, the dynamically generated pole has negative infinite imaginary part, whereas the confinement ground state comes out at $\sqrt{s} = 2.454$ GeV on the real axis. The arrows indicate how the poles move when unquenching increases. The model's physical values are indicated by dots. The imaginary axis is drawn at the DK threshold.

lying poles in the scattering matrix for increasing $c\bar{s}$ - DK coupling [6]. The BABAR collaboration reported in Ref. [21] on the possible existence of a broad $c\bar{s}$ resonance, which might correspond to the dynamically generated pole [22].

Expression (6) thus yields more poles than we bargained for². The jump from Eq. (4) to Eq. (6) contains the physics outlined in Ref. [2]: meson pairs with non-exotic quantum numbers couple to $q\bar{q}$ states through 3P_0 quark-pair annihilation/creation. This mechanism yields the resonances which we expected from the quark-antiquark confinement. But it also yields quasi-bound meson-meson molecules due to shielding caused by the quark-pair annihilation/creation. By model-reducing the intensity of the latter process, the associated poles move into the continuum and disappear from the spectrum of resonances in meson-meson scattering (see Fig. 1). The low-lying scalar mesons belong to this set of resonances [24].

²Recently, related studies have been carried out in Refs. [23]

The two trajectories shown in Fig. 3 come close to each other for certain values of the $c\bar{c}$ - DK coupling. Upon a variation of one other model parameter, this becomes a saddle point. Depending on the value of this parameter, the trajectories may interchange. In that case the end points are connected differently, making the $D_{s0}^*(2317)$ the dynamically generated state, whereas the other pole then seems to stem from the confinement ground state. This is actually what appears to happen for the light positive-parity ground-state mesons and makes them move up in energy when unquenching is turned on. For the scalar mesons, these states correspond to the $f_0(1370)$, $f_0(1500)$, $K_0^*(1430)$, and $a_0(1450)$. The dynamically generated poles correspond to the nonet of lower-lying scalar mesons [25].

5 Conclusions

Most probably, mesons are just mixtures [18, 26] of *quark-antiquark* states, *two-meson molecules*, *glueballs*, *tetraquarks*, *hexaquarks*, *hybrids*, and so forth. We have shown that the first two of the latter list of possible components are the most relevant ones [2, 5–8]. Moreover, a resonance is really a collection of states, all with different masses. Each of these states will have a different composition.

In the spectrum of scattering poles for two-meson systems coupled to $q\bar{q}$ states, we find an infinity of resonances consisting of two distinguishable subsets. One subset manifests the phenomenon of $q\bar{q}$ confinement, whereas the other subset is a direct consequence of quark-pair annihilation/creation. The ground-state scalar nonet of confinement poles is formed by the nonet $f_0(1370)$, $f_0(1500)$, $K_0^*(1430)$, and $a_0(1450)$. On the other hand, the low-lying scalar nonet $f_0(600)$, $f_0(980)$, $K_0^*(800)$, and $a_0(980)$ is the lowest-in-mass scalar nonet of continuum poles.

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